

On the Distinguishability of Relativistic Quantum States in Quantum Cryptography

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Abstract

Relativistic quantum field theory imposes additional fundamental restrictions on the distinguishability of quantum states. Because of the unavoidable delocalization of the quantum field states in the Minkowski space-time, the reliable (with unit probability) distinguishability of orthogonal states formally requires infinite time. For the cryptographic protocols which are finite in time the latter means that the effective “noise” is present even in the ideal communication channel because of the non-localizability of the quantum field states.

PACS numbers: 03.67.-a, 03.65.Bz, 42.50Dv

1 Introduction

Employment of quantum states as the information carriers opened new prospects for the realization of various cryptographic protocols [1–5]. An important result in this field was achieved when it was shown that within the framework of the non-relativistic quantum mechanics some basic classical cryptographic protocols (whose security in classical physics is based on the computational complexity only) become unconditionally secure, i.e. their security is based on the laws of nature (non-relativistic quantum mechanics) and cease to depend on the available computational resources. However, for a number of basic tasks, e.g. bit commitment, all the numerous efforts to develop an unconditionally secure quantum protocol failed [6–9]. Moreover, it was shown later that no ideal unconditionally secure bit commitment protocol can be constructed in the framework of the non-relativistic quantum mechanics [10-11].

The only information carriers suitable for the realistic large-scale quantum cryptographic systems are the photons which are essentially relativistic particles. Since the relativistic quantum mechanics does not allow any sensible physical interpretation, the theory describing the relativistic quantum systems arises as already the quantum field theory [12]. The relativistic quantum field theory imposes additional fundamental restrictions on the speed of information transfer and the processes of the quantum field state measurements. The latter circumstance has the consequence that the secret protocols which cannot be realized in the non-relativistic quantum mechanics become realizable in the classical relativistic theory and in the quantum field theory [13–15].

Below we shall show that in the relativistic quantum field theory the reliable (with unit probability) distinguishability of two orthogonal states requires an infinite time (which actually is a consequence of the microcausality principle), in contrast to the non-relativistic quantum mechanics where this can be done instantly at any time. This circumstance allows to construct new cryptographic protocols in the quantum field theory.

All quantum cryptographic protocols actually employ the following two features of quantum theory. The first one is the no cloning theorem [16], i.e. the impossibility of copying of an arbitrary quantum state which is not known beforehand or, in other words, the impossibility of the following process:

$$|A\rangle|\psi\rangle \rightarrow U(|A\rangle|\psi\rangle) = |B_\psi\rangle|\psi\rangle|\psi\rangle,$$

where $|A\rangle$ and $|B_\psi\rangle$ are the apparatus states before and copying act, respectively, and U is a unitary operator. Such a process is prohibited by the linearity and unitary nature of quantum evolution. Actually, even a weaker process of obtaining any information about one of the two non-orthogonal

states without disturbing it is impossible, i.e. the final states of the apparatus $|A_{\psi_1}\rangle$ and $|A_{\psi_2}\rangle$ corresponding to the initial input states $|\psi_1\rangle$ and $|\psi_2\rangle$, respectively, after the unitary evolution U ,

$$|A\rangle|\psi_1\rangle \rightarrow U(|A\rangle|\psi_1\rangle) = |A_{\psi_1}\rangle|\psi_1\rangle,$$

$$|A\rangle|\psi_2\rangle \rightarrow U(|A\rangle|\psi_2\rangle) = |A_{\psi_2}\rangle|\psi_2\rangle,$$

can only be different, $|A_{\psi_1}\rangle \neq |A_{\psi_2}\rangle$, if $\langle\psi_1|\psi_2\rangle = 0$ [10], which means the impossibility of reliable distinguishing between non-orthogonal states. There is no such a restriction for orthogonal states. Moreover, in the non-relativistic quantum mechanics generally there is no restriction on the instant (arbitrarily fast) reliable distinguishability of any two orthogonal states at any time without disturbing them.

That is why the orthogonal states are not even discussed in the context of non-relativistic quantum cryptographic protocols.

On the contrary, in the quantum field theory it turns out that the distinguishability of two orthogonal states with the probability arbitrarily close to unity requires finite time, and the reliable distinguishability (with unit probability) can only be achieved in infinite time.

To clarify the difference between the non-relativistic quantum mechanics and the relativistic quantum field theory, consider first the non-relativistic case.

Suppose we have a pair of orthogonal states in the Hilbert space: $|\psi_{1,2}\rangle \in \mathcal{H}$ and $\langle\psi_1|\psi_2\rangle = 0$. To reliably distinguish these states, one can for example use the following orthogonal identity resolution in \mathcal{H} :

$$\mathcal{P}_1 + \mathcal{P}_2 + \mathcal{P}_{\perp} = I, \quad \mathcal{P}_{1,2} = |\psi_{1,2}\rangle\langle\psi_{1,2}|, \quad \mathcal{P}_{\perp} = I - \mathcal{P}_1 - \mathcal{P}_2, \quad (1)$$

where $\mathcal{P}_{1,2}$ are the projectors on the subspaces $\mathcal{H}_{1,2}$ spanned by the states $|\psi_1\rangle$ and $|\psi_2\rangle$, while \mathcal{P}_{\perp} is the projector on $\mathcal{H}_{12}^{\perp} = (\mathcal{H}_1 \oplus \mathcal{H}_2)^{\perp}$. The probability of obtaining, for example, the result in the channel \mathcal{P}_1 if in input state was $|\psi_1\rangle$, is

$$\text{Pr}_1\{|\psi_1\rangle\} = \text{Tr}\{|\psi_1\rangle\langle\psi_1|\mathcal{P}_1\} \equiv 1, \quad (2)$$

while it is identically equal to zero in the channels $\mathcal{P}_{2,\perp}$:

$$\text{Pr}_{2,\perp}\{|\psi_1\rangle\} = \text{Tr}\{|\psi_1\rangle\langle\psi_1|\mathcal{P}_{2,\perp}\} \equiv 0, \quad (3)$$

and similarly for the input state $|\psi_2\rangle$. The equations (1–3) mean that the orthogonal states can be reliably (with unit probability) distinguished. To answer the question of whether or not this can be done instantly and without disturbing the measured states, one should consider the measurement procedure in more detail. Indeed, there are three natural levels of description of the measurement process in quantum mechanics which differ in the amount of information they provide [17–19]. The simplest description of the measurement procedure lists only the possible measurement outcomes (i.e., specifies the space of possible measurement outcomes) and specifies the relative frequencies (probabilities) of occurrence of a particular outcome for a given input state of the measured quantum system. In that sense the measurements are in one-to-one correspondence with the positive identity resolutions in the Hilbert space of the system states [17–19]. However, this approach completely ignores the problem of finding the state of the system after the measurement which gave a particular result. At the next level of description of quantum mechanical measurement each measurement procedure is associated with the so-called instrument T which actually is, speaking informally, a rule allowing to ascribe to each pair consisting of the input density matrix ρ_0 of the measured system and the measurement result i a positive matrix $\tilde{\rho}_i$ with the trace $\text{Tr}\tilde{\rho}_i = p_i < 1$, which is interpreted as an unnormalized density matrix of the ensemble of the systems selected by the condition that the performed measurement procedure gave the outcome i (the probability of obtaining the i -th outcome being p_i and the state of the system after the outcome i is obtained being $\rho_i = \tilde{\rho}_i/p_i$). In this approach the von Neumann – Lüders projection postulate is equivalent to the statement that the instrument T corresponding to the measurement described by Eqs.(1–3) is given by the formula

$$\tilde{\rho}_i = \mathcal{P}_i \rho \mathcal{P}_i,$$

so that for example for the input state $|\psi_1\rangle$ we have

$$\tilde{\rho}_1 = \frac{\mathcal{P}_1|\psi_1\rangle\langle\psi_1|\mathcal{P}_1}{\text{Tr}\{|\psi_1\rangle\langle\psi_1|\mathcal{P}_1\}} = |\psi_1\rangle\langle\psi_1|. \quad (4)$$

Equation (4) means that under the assumption of the possibility of the realization of von Neumann measurement satisfying the projection postulate the orthogonal states can be reliably distinguished without any disturbance. One should also note that the reliable distinguishing of orthogonal states without any disturbance can be performed at arbitrarily chosen moment of time, i.e. the measurement procedure can be started at any moment although the procedure itself generally depends on the chosen starting moment. Indeed, since the states $|\psi_1\rangle$ and $|\psi_2\rangle$ are assumed to be known and one should only determine which of these two states is actually given for a particular measurement run, the temporal evolution of these states ($|\psi_i(t)\rangle = U(t)|\psi_i\rangle$, where $|\psi_i\rangle$ is the system state at $t = 0$) should also be considered as a known function of time. If the measurement is performed at time t , the identity resolution should employ $\mathcal{P}_i(t) = U(t)\mathcal{P}_iU^{-1}(t)$ as the corresponding projectors.

An important point is that the duration of the measurement procedure τ has not yet been actually mentioned. To introduce τ , one should consider the third, most detailed, level description of the measurement [18,19]. In the quantum mechanical measurement theory it is shown that any instrument can be represented in the following way: first the measured quantum system S begins to interact with an auxiliary system A (ancilla) and they perform a joint evolution during time τ . Then the system A is subjected to a von Neumann measurement (assumed to be completed arbitrarily fast) whose outcome determines the state of system S immediately after the measurement. Therefore, the speed with which a particular measurement procedure can actually be realized depends on the possibility of choosing a suitable auxiliary system A and the realization of required interaction between A and S .

2 Non-relativistic case

Since the statement of the possibility of reliable distinguishability at any time of the two orthogonal states is quite general, we shall make the transition to the relativistic case smoother by considering in the non-relativistic case two states $|\psi_1\rangle$ and $|\psi_2\rangle$ of a one-dimensional free particle.

It is intuitively clear that to identify a state one should have access to the entire spatial domain where the wavefunction of the particle is present. In the non-relativistic case there are no restrictions on the maximum speed of the information transfer so that an entire extended spatial domain can be accessed instantly at any time, i.e. there are no restrictions on instant non-local measurements.

It is natural to assume that the identification of the state of a quantum field also requires access to the entire domain of the Minkowski space-time where the field is present. However, the existence of the maximum speed of information transfer no extended domain can be accessed instantly, i.e. non-local any measurement requires a finite time.

In the non-relativistic quantum mechanics the above considerations are obvious because the states of a non-relativistic particle can be described by a *wavefunction*. In the relativistic quantum field theory these intuitively appealing considerations should be analyzed more rigorously, since the states of a quantum field cannot be described by a *wavefunction* (or even by an operator field).

The set of observables of a quantum system is represented by an operator algebra with a unit and a conjugation (involution) which are defined on a dense subspace $\Omega \subset \mathcal{H}$ so that $A|\varphi\rangle \in \Omega$ if $|\varphi\rangle \in \Omega$. The Hilbert state \mathcal{H} itself is a completion of Ω with respect to the convergence defined by the norm induced by the scalar product in \mathcal{H} . The convergence conditions in Ω are stronger than in \mathcal{H} and can be chosen in a suitable way for a particular physical system. Continuous functionals defined on elements $|\varphi\rangle \in \Omega \subset \mathcal{H}$ constitute a linear space Ω^* dual to Ω . The space Ω^* contains more elements than \mathcal{H}^* dual to \mathcal{H} . The space \mathcal{H}^* consists of continuous linear functionals defined on elements $|\varphi\rangle \in \mathcal{H}$, frequently written as $\langle f| \in \mathcal{H}^*$ (the value of a functional is a *c*-number $\langle f|\varphi\rangle$). The space \mathcal{H}^* is actually known to be isomorphic to \mathcal{H} itself.

The elements of Ω^* frequently arise as the generalized eigenvectors $|\psi_\alpha\rangle$ of the operator A with a continuous spectrum

$$\begin{aligned} A|\psi_\alpha\rangle &= \lambda_\alpha|\psi_\alpha\rangle, \quad |\psi_\alpha\rangle \in \Omega^*, \\ \langle\psi_\alpha|\psi_{\alpha'}\rangle &= \delta(\alpha - \alpha'). \end{aligned} \quad (5)$$

Any state vector $|\varphi\rangle \in \Omega \subset \mathcal{H}$ can be expanded in a complete set of generalized eigenvectors

$$|\varphi\rangle = \int \varphi(\alpha)|\psi_\alpha\rangle d\alpha, \quad (6)$$

where the coefficients (wavefunctions) $\varphi(\alpha)$ are given by the values of a functional from Ω^*

$$\varphi(\alpha) = \langle\psi_\alpha|\varphi\rangle. \quad (7)$$

The construction $\Omega \subset \mathcal{H} \subset \Omega^*$ is called rigged Hilbert space (Gel'fand triplet) [22–24].

The self-adjoint operator with a continuous spectrum possesses a spectral resolution built of the spectral projectors associated with the generalized eigenvectors

$$A = \int \lambda_\alpha|\psi_\alpha\rangle\langle\psi_\alpha|d\alpha, \quad (8)$$

and the corresponding identity resolution

$$I = \int |\psi_\alpha\rangle\langle\psi_\alpha|d\alpha. \quad (9)$$

A particular functional realization of the rigged space can be chosen in the form of $\mathcal{J}(x) \subset \mathcal{L}_2(x, dx) \subset \mathcal{J}^*(x)$ (where $\mathcal{J}(x)$ is the space of smooth rapidly decreasing functions, $\mathcal{L}_2(x, dx)$ is the space of square integrable functions, and $\mathcal{J}^*(x)$ is the space of tempered distributions [23–25]).

Consider two orthogonal states of a free non-relativistic one-dimensional particle:

$$|\varphi_{1,2}\rangle = \int_{-\infty}^{\infty} \varphi_{1,2}(x)|x\rangle dx, \quad |\varphi_{1,2}\rangle \in \mathcal{J}(x) \subset \mathcal{L}_2(x, dx), \quad (10)$$

$$\langle\varphi_1|\varphi_2\rangle = 0, \quad |x\rangle \in \mathcal{J}^*(x), \quad \langle x|x'\rangle = \delta(x - x'),$$

where $|x\rangle$ is the generalized eigenvector of the position operator. The identity resolution based on the generalized eigenvectors is familiarly written as

$$I = \int_{-\infty}^{\infty} |x\rangle\langle x|dx. \quad (11)$$

The measurement allowing to reliably distinguish between the two states $|\varphi_1\rangle$ and $|\varphi_2\rangle$ without disturbing them is given by the identity resolution defined by Eq.(1). For the measurement outcomes we have

$$\text{Pr}_i\{|\varphi_j\rangle\} = \text{Tr}\{|\varphi_j\rangle\langle\varphi_j|P_j\} = \left| \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \varphi_j^*(x)\delta(x - x')\varphi_i(x')dx dx' \right|^2 = \delta_{ij}. \quad (12)$$

This measurement is non-local in the sense that the reliable state identification requires the access to the entire spatial domain where the wavefunction is different from zero. In the non-relativistic quantum mechanics there is no restriction on the maximum speed of information transfer so that the non-local measurements can be made instantly (arbitrarily fast).

Thus, the physical apparatus implementing the measurement (1–3,12) should be non-local in space. The spatial position of the observer which is the ultimate element in any measurement procedure can be chosen completely arbitrarily since due to the unlimited speed of information transfer the data read off with any non-local device can be instantly gathered to any spatial point where the observed is located.

3 Relativistic case

Consider now the relativistic case. The states of a relativistic quantum system (field) are described by the rays in the physical Hilbert space \mathcal{H} where a unitary representation of the Poincaré group is realized [23,24]. The local quantum field $\varphi(\hat{x})$ (here $\hat{x} = (t, \mathbf{x})$ is a point in the Minkowski space) is defined as a tensor (if the field has more than one component) operator valued distribution. To be more precise, any function (or a set of function if the field has several components) $f(\hat{x}) \in \mathcal{J}(\hat{x})$ is associated with an operator $\varphi(f)$ which can formally be written as

$$\varphi(f) = \sum_{j=1}^r \varphi_j(f_j) = \sum_{j=1}^r \int \varphi_j(\hat{x}) f_j(\hat{x}) d\hat{x}, \quad (13)$$

All the operators $\varphi(f)$ and $\varphi^*(f)$ have a common domain Ω which does not depend on $f(\hat{x})$ and is a dense subspace \mathcal{H} mapped by the operators into itself, $\varphi(f)\Omega \subset \Omega$ ($\varphi^*(f)\Omega \subset \Omega$), so that for the vectors $|\phi\rangle, |\psi\rangle \in \Omega \subset \mathcal{H}$, the quantity $\langle \phi | \varphi(f) | \psi \rangle$ is a distribution from $\mathcal{J}^*(\hat{x})$.

The subspace Ω contains a cyclic vector, called the vacuum state, $|0\rangle \in \Omega$, such that the set of all polynomials $P(\varphi, f)$ (field operator algebra) generate the entire Ω . The field operator algebra elements are defined as

$$P(\varphi, f) = f_0 + \sum_{n=1}^{\infty} \int \int \dots \int \varphi(\hat{x}_1) \varphi(\hat{x}_2) \dots \varphi(\hat{x}_n) f(\hat{x}_1, \hat{x}_2, \dots, \hat{x}_n) d\hat{x}_1 d\hat{x}_2 \dots d\hat{x}_n. \quad (14)$$

The field operators $\varphi(\hat{x})$ map the regular states from Ω to the generalized states $P(\varphi(\hat{x}))\Omega \subset \Omega^*$. The microcausality principle is assumed to be satisfied or, to be more precise, if the functions $f(\hat{x}), g(\hat{y})$ have supports separated by a space-like interval ($\text{supp } f(\hat{x}) \cdot g(\hat{y}) \in (\hat{x} - \hat{y})^2 < 0$), the field operators are assumed to either commute or anticommute, i.e. for any $|\psi\rangle \in \Omega$ the following microcausality relation holds:

$$[\varphi(f), \varphi(g)]_{\pm} |\psi\rangle = 0, \quad (\hat{x} - \hat{y})^2 < 0. \quad (15)$$

At the infinity, the test functions from $\mathcal{J}(\hat{x})$ vanish faster than any polynomial. However, the space of test functions $\mathcal{J}(\hat{x})$ contains a dense set of compactly supported functions $\mathcal{D}(\hat{x}) \subset \mathcal{J}(\hat{x})$ which are zero outside a certain compact domain; therefore, any function from $\mathcal{J}(\hat{x})$ can be approximated by a compactly supported function. The equation (15) should be interpreted as the statement that for the field $\varphi(\hat{x})$ the measurements performed in the domains separated by a space-like interval do not affect each other since no interaction can propagate faster than light.

It is known [23,24] that if one requires that (i) the system states are described by the rays in a Hilbert space where a unitary representation of the Poincaré group is realized and (ii) the spectrum of the group generators lies in the front part of the light cone in the momentum representation then the Lorentz-invariant quantum field can only be realized as an operator-valued distribution rather than the field of operators $\varphi(\hat{x})$ acting in \mathcal{H} . Therefore, one can only make sense of $\varphi(f)$ as an unbounded operator generating a state in $\Omega \subset \mathcal{H}$. The function $f(\hat{x})$ can be interpreted (with some reservations) as the amplitude (“shape”) of a one-particle packet.

The interpretation of a quantized field as a field of operators leads to the trivial two-point function $\langle 0 | \varphi^-(\hat{x}) \varphi^+(\hat{y}) | 0 \rangle = \text{const}$ and violation of the microcausality principle.

In the relativistic case the entire physical Hilbert space \mathcal{H} is a direct sum of the coherent subspaces where different representations of the Poincaré group are realized. Roughly speaking, different subspaces correspond to different types of particles. Considered in the rest of the paper are the one-particle states of a free field; we shall restrict ourselves to the neutral scalar field, spinor field of Dirac electrons and a gauge field (photons).

Suppose we are given two orthogonal states of the neutral scalar field, $|\varphi_{1,2}\rangle \in \Omega \subset \mathcal{H}$

$$|\varphi_{1,2}\rangle = \varphi^+(f_{1,2})|0\rangle = \int \varphi^+(\hat{x}) f_{1,2}(\hat{x}) d\hat{x} |0\rangle, \quad (16)$$

$$\varphi^{\pm}(\hat{x}) = \frac{1}{(2\pi)^{3/2}} \int_{V_m^{\pm}} e^{\pm i \hat{p} \hat{x}} a^{\pm}(\mathbf{p}) \frac{d\mathbf{p}}{\sqrt{2p_0}}, \quad (17)$$

$$|\mathbf{p}\rangle = a^+(\mathbf{p})|0\rangle \in \Omega^*, \quad \varphi^+(\hat{x})|0\rangle \in \Omega^*,$$

where $a^\pm(\mathbf{p})$ are the field operators in the momentum representation which generate the generalized eigenvectors from Ω^* . The integration in Eq.(16) is performed over the front part of the mass shell V_m^+ ($p_0^2 - \mathbf{p}^2 = m^2$, $p_0 > 0$). The operator valued distribution $\varphi^\pm(\hat{x})$ satisfies the Klein-Gordon equation

$$(\square + m^2)\varphi(\hat{x}) = 0, \quad \varphi(\hat{x}) = \varphi^+(\hat{x}) + \varphi^-(\hat{x}). \quad (18)$$

Since in the relativistic case, just in the non-relativistic quantum mechanics, the field states $|\varphi_{1,2}\rangle = \varphi^+(f_{1,2})|0\rangle \in \Omega \subset \mathcal{H}$ are described by the rays in the Hilbert space, the appropriate measurement has the form analogous to Eq.(1) because the orthogonal identity resolution of that kind is only based on the geometrical properties of \mathcal{H} (projection on the rays corresponding to the states $|\varphi_{1,2}\rangle$). The measurement allowing to reliably distinguish between the two orthogonal states is given by the identity resolution

$$\begin{aligned} \mathcal{P}_1 + \mathcal{P}_2 + \mathcal{P}_\perp &= I, \quad \mathcal{P}_\perp = I - \mathcal{P}_1 - \mathcal{P}_2, \\ \mathcal{P}_{1,2} &= |\varphi_{1,2}\rangle\langle\varphi_{1,2}| = \varphi^+(f_{1,2})|0\rangle\langle 0|\varphi^-(f_{1,2}), \\ I &= \int \varphi^+(\hat{x})|0\rangle\langle 0|\varphi^-(\hat{x})d\mathbf{x} = \int_{V_m^+} |\mathbf{p}\rangle\langle\mathbf{p}| \frac{d\mathbf{p}}{2p_0}. \end{aligned} \quad (19)$$

The above identity resolution is written in terms of the generalized eigenstates $\varphi^+(\hat{x})|0\rangle \in \Omega^*$.

The probabilities of obtaining the outcomes in different possible channels are

$$\begin{aligned} \text{Pr}_i\{|\varphi_j\rangle\} &= \text{Tr}\{|\varphi_i\rangle\langle\varphi_i|\mathcal{P}_j\} = |\langle\varphi_j|\varphi_i\rangle|^2 = \left| \int \int f_j^*(\hat{x})D_m^+(\hat{x} - \hat{x}')f_i(\hat{x}')d\hat{x}d\hat{x}' \right|^2 = \\ &= \left| \int f_j^*(\hat{p})f_i(\hat{p})\theta(p_0)\delta(\hat{p}^2 - m^2)d\hat{p} \right|^2 = \left| \int_{V_m^+} f_j^*(\mathbf{p})f_i(\mathbf{p}) \frac{d\mathbf{p}}{2p_0} \right|^2 = \delta_{ij}, \end{aligned} \quad (20)$$

where the generalized commutator function for the field with mass m is [23]

$$\begin{aligned} D_m^\pm(\hat{x}) &= \pm \frac{1}{i(2\pi)^{3/2}} \int e^{i\hat{p}\hat{x}} \theta(\pm p_0)\delta(\hat{p}^2 - m^2)d\hat{p} = \\ \frac{1}{4\pi} \varepsilon(x_0)\delta(\hat{x}^2) &\mp \frac{im}{8\pi\sqrt{\hat{x}^2}} \theta(\hat{x}^2) \left[N_1(m\sqrt{\hat{x}^2}) \mp i\varepsilon(x_0)J_1(m\sqrt{\hat{x}^2}) \right] \pm \frac{im}{4\pi^2\sqrt{-\hat{x}^2}} \theta(-\hat{x}^2)K_1(m\sqrt{-\hat{x}^2}), \\ \varepsilon(x_0)\delta(\hat{x}^2) &\equiv \frac{\delta(x_0 - |\mathbf{x}|) - \delta(x_0 + |\mathbf{x}|)}{2|\mathbf{x}|}. \end{aligned} \quad (21)$$

To within the exponential tails, the commutator function is only different from zero inside the light cone and has a singularity at the cone itself $\lambda^2 = (\hat{x} - \hat{x}')^2 = 0$; outside the light cone the functions $D^\pm(\lambda)$ vanish exponentially at the Compton length as $|\lambda|^{-3/4} \exp(-m\sqrt{|\lambda|})$ [12,23,24]. For a fixed point \hat{x} the contribution to the integral is only given by the points \hat{x}' lying inside the light cone issued from the point \hat{x} , which is actually the consequence of the microcausality principle and reflects the impossibility of the faster-than-light field propagation.

The non-zero tails of the commutator function for the massive particles $m \neq 0$ at the Compton length outside the light cone do not result in any contradictions with the macroscopic causality [26].

The physical states $\varphi^+(f_i)|0\rangle \in \mathcal{H}$ corresponding to two different functions f_i can generally be identical, so that there is no one-to-one correspondence between the functions $f(\hat{x}) \in \mathcal{J}(\hat{x})$ and the states they generate (the generating functions f_i are only recoverable from the state $|\varphi_i\rangle \in \mathcal{H}$ to within the equivalence class). To be more precise, the states are only determined (as it is seen from Eqs.(16,17)) by the values of the generating function in the momentum representation $f(\hat{p})$ on the mass shell $p_0 = \sqrt{\mathbf{p}^2 + m^2}$. The functions $f_{1,2}(\mathbf{p})$ should be interpreted as the restriction of the functions $f_{1,2}(\hat{p})$ to the mass shell, so that the states corresponding to any functions $f(\hat{x})$ coinciding in the \hat{p} -representation on the mass shell are physically identical.

Figure 1:

To eliminate this ambiguity it is convenient to rewrite Eq.(20) in the following equivalent way:

$$\text{Pr}_i\{|\varphi_j\rangle\} = \text{Tr}\{|\varphi_i\rangle\langle\varphi_i|\mathcal{P}_j\} = |\langle\varphi_j|\varphi_i\rangle|^2 = \left| \int f_j^*(\mathbf{x}) D_m^+(\mathbf{x} - \mathbf{x}') f_i(\mathbf{x}') d\mathbf{x} d\mathbf{x}' \right|^2 = \delta_{ij}, \quad (22)$$

where

$$f_i(\mathbf{x}) = \int e^{-i(t_0 p_0(\mathbf{p}) - \mathbf{p}\cdot\mathbf{x})} f_i(\mathbf{p}) d\mathbf{p}. \quad (23)$$

The commutator function $D_m^+(\mathbf{x} - \mathbf{x}')$ is obtained from Eq.(21) if one sets $t = t'$ in $D_m^+(\hat{x} - \hat{x}')$ ($\hat{x} = (t, \mathbf{x})$, $\hat{x}' = (t', \mathbf{x}')$), i.e. the non-zero contributions to the measurement at a given time slice $t = t'$ are only given by spatially coinciding points. Here $f_i(\mathbf{x})$ is understood as the amplitude (playing the role similar to the one-particle wavefunction in the position representation in the non-relativistic case) taken at all points at the same moment of time corresponding to the beginning of the measurement procedure, i.e. to the moment of time starting from which the entire spatial domain embracing the considered state ($f_i(\mathbf{x})$) becomes accessible to the measuring apparatus.

The measurement defined by Eqs.(20,22) is non-local in the sense that it requires the access to the entire domain of the coordinate space \mathbf{x} where the functions $f_i(\mathbf{x})$ are different from zero. Formally, this measurement can be interpreted as a non-local one in the coordinate space performed at a particular moment of time. The device implementing such a measurement should occupy an extended (formally even infinite) domain in the \mathbf{x} -space, i.e. it should have simultaneous access to the entire state. The non-local measurements of that kind are not forbidden in the relativistic case. However, the outcome of the measurement performed with a non-local device cannot be obtained at the same moment as the measurement starts if the measurement encompasses the entire spatial domain where the state is present since the information relevant to the measurement outcome cannot be gathered instantly to the observer located at a certain spatial point from all the points of an extended spatial domain. The information can only be gathered in finite time since actually the indicated spatial domain should be covered by the past part of the light cone issued from the point where the observer is located (Fig.1). The entire domain where the state is present should reside in the interior part of the light cone. The minimum time required for the measurement (for fixed input states) can be determined by examining all the observer positions corresponding to the situation where the entire domain where the state is present at the time t_0 and which is accessible to the measuring apparatus is completely covered by the past part of the light cone.

Since the required measurement time is only determined by the condition of covering the domain by the light cone, it does not depend on the choice of the reference frame because the light cone is a Lorentz-invariant object. The latter can be graphically illustrated for a one-dimensional state. Suppose that the prepared state in a certain reference frame has a characteristic extent $x_0 - x_1 = L$ along the x -axis. The time required in that system is $t = L/2c$ (the optimal observer position is at the center of the domain). At first site, the observer in a moving reference frame would see a Lorentz-contracted domain where the state is defined. Indeed, transformation from the initial reference frame

(x, y, z, t) to the moving one (x', y', z', t') through the hyperbolic rotation (Lorentz transformation)

$$\begin{pmatrix} x' \\ t' \end{pmatrix} = \begin{pmatrix} \operatorname{ch}\psi & \operatorname{sh}\psi \\ \operatorname{sh}\psi & \operatorname{ch}\psi \end{pmatrix} \begin{pmatrix} x \\ t \end{pmatrix}, \quad y' = y, \quad z' = z, \quad \beta = \operatorname{th}\psi, \quad (24)$$

makes the domain size equal to $x'_0 - x'_1 = (x_0 - x_1)\sqrt{1 - \beta^2}$. The time necessary for obtaining the information on the state in the new reference frame is $t'_0 - t'_1 = (x'_0 - x'_1)/2c = (x_0 - x_1)\sqrt{1 - \beta^2}/2c$ and can be done arbitrarily small in the moving reference frame. However, the time elapsed in the initial reference frame $t_0 - t_1 = (t'_0 - t'_1)/\sqrt{1 - \beta^2} = (x_0 - x_1)/2c$ remains the same¹.

For the three-dimensional case the minimal required time will be determined by the size of the maximum cross-section of the spatial domain.

The functions $f_i(\mathbf{x})$ rapidly decrease at the infinity, but do not become identical zero outside of any compact domain actually because it is impossible to obtain a function with compact support in the position space taking a Fourier transform of a function which is only defined on the mass shell in the momentum space (the proof based on the Wiener–Paley theorem can be found, e.g. in Refs.[27,28]). Formally, this means that the time necessary for the reliable identification (with unit probability) of one of the two orthogonal field states is infinite since because of the infinite support of the functions generating the field states one should have access to the entire coordinate space. However, the states still can be identified in a finite time (which, of course, depends on the structure of the chosen states) with the probability arbitrarily close to unity.

Consider now the case of a multicomponent spinor field of Dirac electrons. The operator valued spinor field distribution has the form [23,24]

$$\psi(\hat{x}) = \frac{\sqrt{m}}{(2\pi)^{3/2}} \int \frac{d\mathbf{p}}{\sqrt{\varepsilon_{\mathbf{p}}}} \sum_{\zeta=\pm 1/2} \left\{ a_{\zeta}^{(+)}(\mathbf{p}) u_{\zeta}^{(+)}(\mathbf{p}) e^{-i\hat{p}\hat{x}} + a_{\zeta}^{*(-)}(\mathbf{p}) u_{\zeta}^{(-)}(\mathbf{p}) e^{i\hat{p}\hat{x}} \right\}, \quad (25)$$

and the Dirac conjugate operator is

$$\tilde{\psi}(\hat{x}) = \frac{\sqrt{m}}{(2\pi)^{3/2}} \int \frac{d\mathbf{p}}{\sqrt{\varepsilon_{\mathbf{p}}}} \sum_{\zeta=\pm 1/2} \left\{ a_{\zeta}^{*(+)}(\mathbf{p}) \tilde{u}_{\zeta}^{(+)}(\mathbf{p}) e^{i\hat{p}\hat{x}} + a_{\zeta}^{(-)}(\mathbf{p}) \tilde{u}_{\zeta}^{(-)}(\mathbf{p}) e^{-i\hat{p}\hat{x}} \right\} \quad (26)$$

with the normalization conditions of the Dirac spinors

$$2 \sum_{\zeta=\pm 1/2} u_{\zeta}^{(\pm)\alpha}(\mathbf{p}) \tilde{u}_{\zeta\beta}^{(\pm)}(\mathbf{p}) = \left(\frac{\hat{p}}{m} \right)_{\beta}^{\alpha} \pm \delta_{\beta}^{\alpha}, \quad \varepsilon_{\mathbf{p}} = \sqrt{\mathbf{p}^2 + m^2}, \quad \alpha, \beta = 1 \dots 4. \quad (27)$$

The operator valued distribution (25) satisfies the Dirac equation

$$(\hat{p} + m)\psi(\hat{x}) = 0, \quad \hat{p} = i\gamma^{\mu}\partial_{\mu}, \quad (28)$$

and the anticommutation relations

$$[\psi^{\alpha}(\hat{x}), \psi^{\beta}(\hat{x}')]_{+} = [\tilde{\psi}_{\alpha}(\hat{x}), \tilde{\psi}_{\beta}(\hat{x}')]_{+} = 0,$$

$$[\psi^{\alpha}(\hat{x}), \tilde{\psi}_{\beta}(\hat{x}')]_{+} = -iS_{\beta}^{\alpha}(\hat{x} - \hat{x}') = (i\gamma^{\mu}\partial_{\mu} + m)D_m(\hat{x} - \hat{x}'), \quad D_m(\hat{x}) = D_m^{+}(\hat{x}) + D_m^{-}(\hat{x}). \quad (29)$$

The smeared operator functions can be written in the form [12]

$$|\psi(\mathbf{f})\rangle = \int \psi^{\alpha}(\hat{x}) f^{\alpha}(\hat{x}) d\hat{x}, \quad f^{\alpha}(\hat{x}) \in \mathcal{J}(\hat{x}). \quad (30)$$

The two orthogonal states $|\varphi_{1,2}\rangle$ ($\langle\varphi_1|\varphi_2\rangle = 0$) of the spinor field can be written

$$|\varphi_{1,2}\rangle = \left(\int \psi^{\alpha}(\hat{x}) f^{\alpha}(\hat{x}) d\hat{x} \right) |0\rangle \in \Omega \subset \mathcal{H}, \quad (31)$$

¹This point is important for construction of cryptographic protocols, for example, for the quantum cryptography based on orthogonal states, since it prohibits eavesdropping employing the twin paradox.

with the corresponding identity resolution employing the generalized states

$$I = \int (\psi(\hat{x})) |0\rangle\langle 0| (\tilde{\psi}(\hat{x})) d\mathbf{x} = \sum_{\zeta=\pm\frac{1}{2}, s=\pm} \int |\mathbf{p}, \zeta, s\rangle\langle \mathbf{p}, \zeta, s| \frac{md\mathbf{p}}{2\varepsilon(\mathbf{p})}, \quad (32)$$

$$|\mathbf{p}, \zeta, s\rangle = a_\zeta^{*(s)}(\mathbf{p})|0\rangle,$$

and similarly for the projection operators on the states $|\varphi_1\rangle$ and $|\varphi_2\rangle$.

The probabilities of obtaining different measurement outcomes are

$$\begin{aligned} \text{Pr}_i\{|\varphi_j\rangle\} &= \text{Tr}\{|\varphi_i\rangle\langle\varphi_i|\mathcal{P}_j\} = \left| \int \int f_j^{*\alpha}(\hat{x}) S_{m\beta}^{+\alpha}(\hat{x} - \hat{x}') f_{i\beta}(\hat{x}') d\hat{x} d\hat{x}' \right|^2 = \\ &= \left| \int f_j^{*\alpha}(\hat{p})(\hat{p} - m)^\alpha_\beta f_{i\beta}(\hat{p}) \theta(p_0) \delta(\hat{p}^2 - m^2) d\hat{p} \right|^2 = \left| \int_{V_m^+} f_j^{*\alpha}(\mathbf{p})(\hat{p} - m)^\alpha_\beta f_{i\beta}(\mathbf{p}) \frac{d\mathbf{p}}{2\varepsilon(\mathbf{p})} \right|^2 = \delta_{ij}, \end{aligned} \quad (33)$$

Here $f_{i\beta,j\alpha}(\mathbf{p})$ are the values of the amplitude on the mass shell with positive energy. In the present case, just as for the scalar field, the answer is expressed through the derivative of the commutator function $D^\pm(\hat{x} - \hat{x}')$ so that everything said above on the finiteness of the time required for the distinguishing between two states applies as well to the fermionic field.

Consider now the case of a gauge field which is most interesting from the viewpoint of applications, i.e. the photon field. The electromagnetic field operators are written as [12]

$$A_\mu^\pm(\hat{x}) = \frac{1}{(2\pi)^{3/2}} \int \frac{d\mathbf{k}}{\sqrt{2k_0}} e^{\pm i\hat{k}\hat{x}} e_\mu^m(\mathbf{k}) a_m^\pm(\mathbf{k}) \quad (34)$$

and satisfy the commutation relations

$$[A_\mu^-(\hat{x}), A_\nu^-(\hat{x}')]_- = ig_{\mu\nu} D_0^-(\hat{x} - \hat{x}'), \quad (35)$$

where $D_0^-(\hat{x} - \hat{x}')$ is the commutator function for the massless field (21). There are four types of photons: two transverse, one longitudinal, and one temporal. The two latter types are actually fictitious particles and can be eliminated at the expense of introducing an indefinite metric [12]. For our purposes the shortest way to the required result consists in employing a specific gauge. We shall further work in the subspace of physical space using the Coulomb gauge $A_\mu = (\mathbf{A}, \varphi = 0)$ dealing with the two physical transverse states of the electromagnetic field.

The operator valued distribution is a vector in the three-dimensional space:

$$\vec{\psi}(\hat{x}) = \frac{1}{(2\pi)^{3/2}} \int_{V_0^+} \frac{d\mathbf{k}}{\sqrt{2k_0}} \sum_{s=\pm 1} \mathbf{w}(\mathbf{k}, s) \{a(\mathbf{k}, s) e^{-i\hat{k}\hat{x}} + a^\dagger(\mathbf{k}, -s) e^{i\hat{k}\hat{x}}\}; \quad (36)$$

here $\mathbf{w}(\mathbf{k}, s)$ is a three-dimensional vector describing the polarization state $s = \pm 1$,

$$\mathbf{w}(\mathbf{k}, \pm) = \frac{1}{\sqrt{2}} [\mathbf{e}_1(\mathbf{k}) \pm i\mathbf{e}_2(\mathbf{k})], \quad \mathbf{e}_1(\mathbf{k}) \perp \mathbf{e}_2(\mathbf{k}), \quad |\mathbf{w}(\mathbf{k}, s)|^2 = 1, \quad (37)$$

where $\mathbf{e}_{1,2}(\mathbf{k})$ are the orthogonal vectors normal to \mathbf{k} . The field operator satisfies the Maxwell equations

$$\nabla \times \vec{\psi}(\hat{x}) = -i \frac{\partial}{\partial t} \vec{\psi}(\hat{x}), \quad (38)$$

$$\nabla \cdot \vec{\psi}(\hat{x}) = 0.$$

The smeared field operators can be written as

$$\vec{\psi}(f_{1,2}) = \sum_{s=\pm 1} \int \vec{\psi}(\hat{x}, s) f_{1,2}(\hat{x}, s) d\hat{x} = \quad (39)$$

$$\frac{1}{(2\pi)^{3/2}} \int_{V_0^+} \frac{d\mathbf{k}}{\sqrt{2k_0}} \sum_{s=\pm 1} \mathbf{w}(\mathbf{k}, s) \{ f_{1,2}(\mathbf{k}, s) a(\mathbf{k}, s) e^{-i\hat{k}\hat{x}} + f_{1,2}(\mathbf{k}, -s) a^+(\mathbf{k}, -s) e^{i\hat{k}\hat{x}} \},$$

where the values of the functions $f_{1,2}(\mathbf{k}, s)$ are taken at the mass shell (light cone V_0^+ in the momentum representation). Two orthogonal states of the photon field can be written in the form

$$|\vec{\psi}_{1,2}\rangle = (\vec{\psi}(f_{1,2})) |0\rangle. \quad (40)$$

The corresponding orthogonal projectors on the states $|\vec{\psi}_{1,2}\rangle$ are

$$\mathcal{P}_{1,2} = |\vec{\psi}_{1,2}\rangle \langle \vec{\psi}_{1,2}|, \quad (41)$$

$$I = \int (\vec{\psi}^+(\hat{x})) |0\rangle \langle 0| (\vec{\psi}^-(\hat{x})) d\mathbf{x} = \sum_{s=\pm 1} \int_{V_0^+} (\mathbf{w}(\mathbf{k}, s) a^+(\mathbf{k}, s)) |0\rangle \langle 0| (\mathbf{w}(\mathbf{k}, s) a(\mathbf{k}, s)) \frac{d\mathbf{k}}{2|\mathbf{k}|}. \quad (42)$$

The probabilities of obtaining different measurement outcomes are

$$\text{Pr}_i\{|\vec{\psi}_j\rangle\} = \text{Tr}\{|\vec{\psi}_i\rangle \langle \vec{\psi}_i| \mathcal{P}_j\} = |\langle \vec{\psi}_j | \vec{\psi}_i \rangle|^2 = \quad (43)$$

$$|\int \int f_j^*(\mathbf{x}) D_0^+(\mathbf{x} - \mathbf{x}') f_i(\mathbf{x}') d\mathbf{x} d\mathbf{x}'|^2 = |\int_{V_0^+} f_i^*(\mathbf{k}) f_j(\mathbf{k}) \frac{d\mathbf{k}}{2|\mathbf{k}|}|^2 = \delta_{ij},$$

where $D_0^+(\mathbf{x} - \mathbf{x}')$ is the commutator function for the massless field at the time slice $t = t'$ has the form

$$D_0^+(\mathbf{x} - \mathbf{x}') = -\frac{1}{4\pi} \frac{\delta(|\mathbf{x} - \mathbf{x}'|)}{2|\mathbf{x} - \mathbf{x}'|}.$$

It should be noted that arising in the non-relativistic case instead of the commutator function D_0^+ is the usual $\delta(\mathbf{x} - \mathbf{x}')$ -function (see Eq.(12)). The latter is actually related to the fact that in the non-relativistic case the integration in the scalar product in the momentum representation is performed with the Galilean-invariant measure $d\mu(\mathbf{p}) = d\mathbf{p}$ instead of the Lorentz-invariant measure $d\mu(\hat{p}) = \theta(p_0)\delta(\hat{p}^2)d\hat{p}$ residing at the mass shell in the relativistic case.

The time necessary for the reliable identification of one of the pair of orthogonal states is determined by the spatial localization of the generating functions $f_{1,2}(\mathbf{k})$ in the coordinate (position) representation or, to be more precise, by the localization of $|\mathbf{f}_{1,2}(\mathbf{x}, t, s)|^2$ $|\mathbf{f}_{1,2}(\mathbf{x}, t, s)|$

$$\mathbf{f}_{1,2}(\mathbf{x}, t, s) = \int_{V_0^+} \mathbf{w}(\mathbf{k}, s) f_{1,2}(\mathbf{k}, s) e^{-i(|\mathbf{k}|t - \mathbf{k}\mathbf{x})} d\mathbf{k}, \quad (44)$$

where the integration is performed over the mass shell V_0^+ (surface of the light cone in the momentum representation). It has long been known (e.g., see Ref.[29] and references therein) that it is impossible to obtain a strictly localizable (with compact support) function in the coordinate space from a normalized function defined on the mass shell. The latter holds for both massive and massless particles. However, there exist the functions with arbitrarily close to the exponential fall off at the infinity. Recently, the one-particle photon states of that kind were explicitly constructed [27]. The restrictions on the arbitrarily close to the exponential fall off stem from the Wiener-Paley theorem for square integrable functions [29] and the impossibility of mixing of the positive and negative frequency states on the mass shell for one-particle states.

The possibility of the existence of the states whose localization is arbitrarily close to the exponential one is actually related to the choice of the test function space $\mathcal{J}(\hat{x})$ which contains the dense subspace of compactly supported functions $\mathcal{D}(\hat{x})$. The latter circumstance is closely related to the local nature of the theory since the existence of a dense set of compactly supported functions defined in the Minkowski \hat{x} -space allows to achieve the local properties of the distributions, including the commutator functions, appearing in the microcausality principle.

Reduction of the set of the test functions can result in the non-local nature of the theory (e.g., see Refs.[23,29–32]). However, non-locality of the theory does not imply violation of the causality principle at the macroscopic level (at the level of the observer) [29,32].

4 Conclusion

Thus, the reliable (with unit probability) identification of one-particle field states generally requires an infinite time because of the requirement of having access to the entire spatial domain where the state “is present”. Formally, the construction of the measurement at the level of the corresponding identity resolution in the relativistic case is completely analogous to the non-relativistic quantum mechanics since in both cases the states are described by the rays in the Hilbert state. At this level, the description of measurement involves only the geometrical properties of the state space (the scalar product and projections on the corresponding orthogonal states). The difference between the non-relativistic and relativistic cases arises at the level of the internal structure of the scalar product and conveying the measurement outcome to the observer. In the non-relativistic case, there are no restrictions on the instantaneous spatially non-local measurements and conveying their outcomes to the observer (which can be located at an arbitrary point) since there are no limits on the maximum speed of information transfer. In the relativistic case the Lorentz-invariant scalar product already contains (through the commutator functions reflecting the microcausality principle) the information on the maximum speed of the field propagation. The latter should only be interpreted as the impossibility for the field to propagate into an extended device (fill it) faster than light. A similar situation occurs if an extended device moves into the domain where the field is present. However, the presence of the field in the device alone does not yet mean obtaining of the measurement outcome. The measurement outcome should be conveyed to a single point (location of the observer). The latter can only be done in a finite time depending on the spatial localization of the states since the spatial domain should be “covered” by the past part of the light cone issued from the point where the observer is situated (the observer can be interpreted as a classical device registering the final measurement outcome).

In the above approach all the information on the non-local (extended) device is contained in the projectors on the corresponding rays in the Hilbert space. Formally, these projectors should be considered as an (non-local) observable which, from the viewpoint of the measurement theory, is quite similar to a non-local in the coordinate space observable (e.g., momentum). Therefore, associated with this observable should be a physical apparatus implementing the corresponding measurement.

The conclusion on the finiteness of the time interval required for the state identification does not depend on a particular measurement procedure. It is often experimentally easier to realize a non-local measurement as a local one employing a localized auxiliary system A interacting with the measured system (it is known that any measurement can be realized in this way [18,19]). In the course of the joint evolution the state of the auxiliary system is gradually changed and finally the measurement is performed over that system. For the orthogonal states there are no restrictions on the existence of non-disturbing measurements [3] (see Eqs.(1–3)) in the sense that just after the measurement only the state of the auxiliary system A is changed while state of the studied system S coincides with its initial state just before the interaction between S and A was turned on. In that case it is also intuitively clear that a finite time is required (if the field state is extended) for the field to pass through a point where the system A is localized and the interaction between A and S occurs which cannot be done faster than $t \sim L/c$ because of the existence of a maximum speed of the field propagation. It is also clear that this time cannot be reduced by enhancement of the local coupling between A and S because this time is limited by the field propagation speed. The required time cannot also be reduced employing the Lorentz contraction of the domain where the field is present occurring in a moving reference frame (see discussion above). Naturally, the auxiliary system A should also be described by the quantum field theory and one should assume that it is so strongly spatially localized that the time necessary to perform a measurement on this system can be done arbitrarily small. However, it is much more difficult to obtain any general results in this approach since it inevitably requires making certain assumptions concerning the system S itself and its interaction with the auxiliary system A .

Formally, the von Neumann measurement described by the orthogonal resolution of identity given by Eq.(43) in the subspace of one-particle states in the Hilbert space does not disturb the orthogonal states since the situation here is analogous to the non-relativistic case described by Eqs.(1–4). However, because of the non-local nature of the projectors, the reliable and non-disturbing identification of the

orthogonal states achieved by this measurement formally requires infinite time.

With respect to the time necessary for the measurement, the situation is quite similar for the non-orthogonal states, the only difference being that the latter cannot be reliably distinguished even in an infinite time.

In conclusion, it should be noted that the non-relativistic quantum cryptographic protocols implicitly assume the possibility of performing a measurement distinguishing between the two states in a finite time (ideally, an instantaneous measurement). Briefly, all the non-relativistic protocols can be described in the following way. User A sends either the state $|\psi_1\rangle$ or $|\psi_2\rangle$ to user B who performs quantum-mechanical measurements. Part of information on the measurement outcomes is discussed through the open communication channel. It is assumed that the states $|\psi_1\rangle$ or $|\psi_2\rangle$ are prepared and measured in certain agreed moments of time so that the preparation and measurement procedures can be performed at arbitrary moments of time. Thus, it is implicitly assumed that the protocol can be implemented in a finite time interval.

The relativistic quantum field theory does not allow access to the entire state in any finite time because of the non-localizable nature of the state. Therefore, all the exchange protocols with finite duration should inevitably involve an error in the preparation and detection of both orthogonal and non-orthogonal one-particle states even for the ideal communication channel. In other words, the non-localizability of the states results in the effective “noise” because of the impossibility of the reliable detection of orthogonal states in a finite time although for a fixed measurement duration the error can be made arbitrarily (exponentially) small by choosing more and more strongly localized state. This circumstance should be taken into account in the realistic cryptographic protocols employing photons as the information carriers.

The authors are grateful to W.L.Golo, S.V.Iordanskii and V.M.Edelstein for the interest and discussion of the results of this work.

This work was supported by the Russian Foundation for Basic Research (project No 99-02-18127), the project “Physical Principles of the Quantum Computer”, and the program “Advanced Devices and Technologies in Micro- and Nanoelectronics” (project No 02.04.5.2.40.T.50).

This work was also supported by the Wihuri Foundation, Finland.

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